EFFECT OF BINDING IN DEEP INELASTIC SCATTERING REVISITED

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Abstract

In the Bjorken limit of the present theory of deep inelastic scattering (DIS) the structure functions (up to anomalous dimensions and perturbative QCD corrections) are described by the parton model. However the current operator in the parton model does not properly commute with the representation operators corresponding to the Lorentz group, space reflection and time reversal. To investigate the violation of these symmetries in the parton model we consider a model in which the current operator explicitly satisfies extended Poincare invariance and current conservation. It is shown that due to binding of quarks in the nucleon the Bjorken variable x no longer can be interpreted as the internal light cone momentum fraction ξ even in the Bjorken limit. As a result, the data on DIS alone do not make it possible to determine the ξ distribution of quarks in the nucleon. We also consider a qualitative explanation of the fact that in the parton model the values given by the sum rules exceed the corresponding experimental quantities while the quark contribution to the nucleon momentum and spin is underestimated.

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1 The statement of the problem

A full description of any relativistic quantum system implies in particular that one can construct the momentum and angular momentum operators $(\hat{P}^{\mu}, \hat{M}^{\mu\nu})$ $(\mu, \nu = 0, 1, 2, 3,\ \hat{M}^{\nu\mu} = -\hat{M}^{\mu\nu})$ which are the generators of the (pseudo)unitary representation of the Poincare group for this system. We

shall always assume that the commutation relations between the generators are realized in the form

$$[\hat{P}^\mu, \hat{P}^\nu] = 0, \quad [\hat{M}^{\mu\nu}, \hat{P}^\rho] = -\imath(\eta^{\mu\rho}\hat{P}^\nu - \eta^{\nu\rho}\hat{P}^\mu),$$

$$[\hat{M}^{\mu\nu}, \hat{M}^{\rho\sigma}] = -\imath(\eta^{\mu\rho}\hat{M}^{\nu\sigma} + \eta^{\nu\sigma}\hat{M}^{\mu\rho} - \eta^{\mu\sigma}\hat{M}^{\nu\rho} - \eta^{\nu\rho}\hat{M}^{\mu\sigma}) \tag{1}$$

where the metric tensor in Minkowski space has the nonzero components $\eta^{00} = -\eta^{11} = -\eta^{22} = -\eta^{33} = 1$ and we use the system of units with $\hbar = c = 1$. We use $(P^\mu, M^{\mu\nu})$ to denote the corresponding generators in the case when all interactions in the system under consideration are turned off.

In the relativistic theory, in contrast with the nonrelativistic one, it is not possible to realize the operators $(\hat{P}^\mu, \hat{M}^{\mu\nu})$ in such a way that only the Hamiltonian $\hat{P}^0$ is interaction dependent while all the other nine generators are free. Indeed, suppose that $\hat{P}^0$ is interaction dependent and consider the relation $[\hat{M}^{0i}, \hat{P}^k] = -\imath\delta_{ik}\hat{P}^0$ $(i, k = 1, 2, 3)$ which follows from Eq. (1). Then it is obvious that if all the operators $\hat{P}^k$ are free then all the operators $\hat{M}^{0i}$ are inevitably interaction dependent and *vice versa*, if all the operators $\hat{M}^{0i}$ are free then all the operators $\hat{P}^k$ are inevitably interaction dependent. According to the Dirac classification [1], the realization of the algebra (1) in such a way that the operators $(\hat{P}^0, \hat{M}^{0i})$ are interaction dependent and the other six generators of the Poincare group are free is called the instant form of dynamics, while the point form implies that all the components $\hat{P}^\mu$ are interaction dependent and $\hat{M}^{\mu\nu} = M^{\mu\nu}$. Instead of the $0, z$ components of four-vectors we also can work with the $\pm$ components defined as $p^\pm = (p^0 \pm p^z)/\sqrt{2}$. Suppose that $\hat{P}^-$ is interaction dependent and the other components of $\hat{P}^\mu$ are free. Then, as follows from Eq. (1), $[\hat{M}^{-j}, \hat{P}^l] = -\imath\delta_{jl}\hat{P}^-$ $(j = 1, 2)$ and hence all the operators $\hat{M}^{-j}$ are inevitably interaction dependent. The realization of the algebra (1) in such a way that $(\hat{P}^-, \hat{M}^{-j})$ are interaction dependent and the other seven generators are free is called the front form [1]. Of course, physical results should not depend on the choice of the form of dynamics and in the general case all the ten generators can be interaction dependent.

We can also consider extended relativistic invariance which implies that the generators properly commute not only with each other but also with the representation operators $\hat{U}_P$ and $\hat{U}_T$ corresponding to space reflection and time reversal. A possible choice in the instant and point forms is $\hat{U}_P = U_P$ and $\hat{U}_T = U_T$, but in the front form the operators $\hat{U}_P$ and $\hat{U}_T$ should

be necessarily interaction dependent. This follows in particular from the relations

$$\hat{U}_P P^+ \hat{U}_P^{-1} = \hat{U}_T P^+ \hat{U}_T^{-1} = \hat{P}^- \tag{2}$$

As noted by Coester [2], the interaction dependence of the operators $\hat{U}_P$ and $\hat{U}_T$ in the front form does not mean that there are discrete dynamical symmetries in addition to the rotations about transverse axes. Indeed, the discrete transformation P_2 such that $P_2\,x := \{x^0, x_1, -x_2, x_3\}$ leaves the light front $x^+ = 0$ invariant and therefore P_2 can be chosen free. The full space reflection P is the product of P_2 and a rotation about the 2-axis by π. Thus it is not an independent dynamical transformation in addition to the rotations about transverse axes. Similarly the transformation TP leaves $x^+ = 0$ invariant and $T = (TP)P_2R_2(\pi)$.

To describe some process of deep inelastic scattering (DIS) it is necessary to construct the operator of the electromagnetic or weak current responsible for the transition *nucleon* $\rightarrow$ *hadrons* in this process. Let $\hat{J}^\mu(x)$ be such an operator where x is a point in Minkowski space. Translational invariance of the current operator implies that

$$exp(-\imath \hat{P}_\mu a^\mu)\hat{J}^\mu(x) exp(\imath \hat{P}_\mu a^\mu) = \hat{J}^\mu(x-a) \tag{3}$$

and Lorentz invariance implies that

$$[\hat{M}^{\mu\nu}, \hat{J}^\rho(x)] = -\imath\{(x^\mu \partial^\nu - x^\nu \partial^\mu)\hat{J}^\rho(x) + \eta^{\mu\rho}\hat{J}^\nu(x) - \eta^{\nu\rho}\hat{J}^\mu(x)\} \tag{4}$$

The electromagnetic current operator should also satisfy the continuity equation $\partial \hat{J}^\mu(x)/\partial x^\mu = 0$. As follows from Eq. (3), this equation can be written in the form

$$[\hat{J}^\mu(x), \hat{P}_\mu] = 0 \tag{5}$$

Finally, the usual requirement is that the theory should be local in the sense that the commutator $[\hat{J}^\mu(x), \hat{J}^\nu(y)]$ should necessarily vanish when $x-y$ is a space-like vector. Let us note however that if a theory is nonlocal in the above sense, this does not necessarily imply that it is nonphysical. Indeed, as it has become clear already in 30th, in relativistic quantum theory there is no operator possessing all the properties of the position operator. In particular, the quantity x in the Lagrangian density $L(x)$ is not the coordinate, but some parameter which becomes the coordinate only in the classical limit. Therefore the physical condition is that only the macroscopic locality should

be satisfied, i.e. the above commutator should vanish when $|(x-y)^2| \to \infty$ but $(x-y)^2 < 0$.

In the framework of the scattering theory it is sufficient to ensure relativistic invariance of the S-matrix and hence it is sufficient to consider Eqs. (1-5) only in the scattering space of the system under consideration. In QED the electrons, positrons and photons are the fundamental particles, and the scattering space is the space of these almost free particles (”in” or ”out” space). Therefore it is sufficient to deal only with $\hat{P}^{\mu}_{ex}, \hat{M}^{\mu\nu}_{ex}$ where ”ex” stands either for ”in” or ”out”. However in QCD the scattering space by no means can be considered as a space of almost free fundamental particles — quarks and gluons. For example, even if the scattering space consists of one particle (say the nucleon), this particle is the bound state of quarks and gluons, and the operators $\hat{P}^{\mu}$ and $\hat{M}^{\mu\nu}$ considerably differ from P^{μ} and $M^{\mu\nu}$. It is well-known that perturbation theory does not apply to bound states and therefore $\hat{P}^{\mu}$ and $\hat{M}^{\mu\nu}$ cannot be determined in the framework of perturbation theory. For these reasons we will be interested in cases when the representation operators correspond to the full generators $\hat{P}^{\mu}$ and $\hat{M}^{\mu\nu}$.

In quantum field theory the only known way of constructing the operators $\hat{P}^{\mu}$, $\hat{M}^{\mu\nu}$ and $\hat{J}^{\mu}(x)$ is canonical formalism. However this formalism ignores the fact that local field operators are operator valued distributions and therefore products of these operators at coinciding points are not well defined. Moreover, canonical formalism contradicts the Haag theorem [3]. A detailed discussion of these problems can be found, for example, in the well-known monographs [4]. It has been also shown in a vast literature (see e.g. Refs. [5, 6, 7, 8, 9]) that the canonical treatment of equal time commutation relations between the current operators can often lead to incorrect results. The results of the above references show that it is often premature to trust assumptions which may seem physical or natural.

In Sect. 2 we argue that the assumptions used in the present theory of DIS are not substantiated and the current operator contains a nontrivial nonperturbative part which contributes to DIS even in the Bjorken limit. Since this contribution cannot be determined at the present stage of QCD, it is reasonable to consider models in which the operators $\hat{P}^{\mu}$, $\hat{M}^{\mu\nu}$ and $\hat{J}^{\mu}(x)$ are well defined and can be explicitly constructed but of course one should understand that such models cannot have any pretensions to be fundamental. In Sect. 3 the explicit construction of the operators $\hat{P}^{\mu}$ and $\hat{M}^{\mu\nu}$ in the framework of our models is considered. In Sect. 4 the well-known results

that the parton model is a consequence of impulse approximation (IA) for the current operator in the front form are reproduced but in our consideration it is clear that such a current operator does not properly commute with the Lorentz group generators and the operators corresponding to the discrete symmetries. As shown in Sect. 5, the idea that at large momentum transfer the current operator can be taken in IA can be consistently realized in the point form and the corresponding results considerably differ from the results of the parton model. In Sect. 6 the difference between the parton model sum rules and the sum rules in our model is considered and Sect. 7 is discussion.

2 Problems in the present theory of DIS

In addition to the examples considered in Refs. [6, 7, 8, 9] we now give one more example which shows that the reader thinking that it is not reasonable to worry about mathematical rigor will be confronted with the following contradiction.

In QED and QCD the electromagnetic current operator is usually written in the form $\hat{J}^\mu(x) = \mathcal{N}\{\hat{\bar{\psi}}(x)\gamma^\mu\hat{\psi}(x)\}$ where $\mathcal{N}$ stands for the normal product, $\hat{\psi}(x)$ is the Heisenberg operator of the Dirac field and for simplicity we do not write flavor operators and color and flavor indices. As noted above, such a definition ignores the fact that $\hat{J}^\mu(x)$ should be an operator valued distribution. Then we can consider $\hat{J}^\mu(x)$ at fixed values of x, for example at $x = 0$, and formally it follows from Eq. (4) that

$$[\hat{M}^{\mu\nu}, \hat{J}^\rho(0)] = -\imath(\eta^{\mu\rho}\hat{J}^\nu(0) - \eta^{\nu\rho}\hat{J}^\mu(0)) \tag{6}$$

Let us now take into account that in the framework of canonical quantization the Heisenberg and Schrodinger pictures at $x = 0$ are the same and therefore $\hat{J}^\mu(0)$ does not depend on the interaction, i.e. $\hat{J}^\mu(0) = J^\mu(0)$ (note that the free operator $J^\mu(x)$ can be considered as a usual operator valued function of x). The operator $\hat{M}^{\mu\nu}$ can be written as $M^{\mu\nu} + V^{\mu\nu}$. Since $[M^{\mu\nu}, J^\rho(0)]$ can be written by analogy with Eq. (6), it is obvious that $[V^{\mu\nu}, J^\rho(0)] = 0$ and in particular $[V^{0i}, J^0(0)] = 0$. The canonical form of V^{0i} in QED is [10]

$$V^{0i} = e\int x^i \mathbf{A}(\mathbf{x})\mathbf{J}(\mathbf{x})d^3\mathbf{x} \tag{7}$$

where e is the electric charge, $\mathbf{A}$ is the operator of the Maxwell field and $\mathbf{x}$ is the spatial part of the four-vector x when $x^0 = 0$. Therefore $[V^{0i}, J^0(0)] = 0$ if

$$\int x^i \mathbf{A}(\mathbf{x})[\mathbf{J}(\mathbf{x}), J^0(0)] d^3\mathbf{x} = 0 \tag{8}$$

It is well-known that if the standard equal-time commutation relations are used naively then the commutator in Eq. (8) vanishes and therefore this equation is satisfied. However when $\mathbf{x} \to 0$ this commutator involves the product of four Dirac fields at $\mathbf{x} = 0$. The famous Schwinger result [5] is that

$$[J^i(\mathbf{x}), J^0(0)] = C \frac{\partial}{\partial x^i} \delta(\mathbf{x}) \tag{9}$$

where C is some (infinite) constant. Therefore Eq. (8) is not satisfied and the current operator $\hat{J}^\mu(x)$ constructed in the framework of canonical formalism does not satisfy Lorentz invariance. As a result, the operator $\hat{J}^\mu(0)$ (if it exists) cannot be free.

Although our example is the demonstration of this property in spinor QED, as noted above, here this example is not very important. However our example is important for QCD since it shows that the algebraic reasons based on Eq. (6) are more solid than the reasons based on formal manipulations with local operators. Indeed, if the operators $\hat{M}^{\mu\nu}$ are interaction dependent then $\hat{J}^\mu(0)$ cannot be free simply because there is no reason for interaction terms in $\hat{M}^{\mu\nu}$ to commute with all the operators $J^\mu(0)$. Therefore in the instant and front forms $\hat{J}^\mu(0) \neq J^\mu(0)$. At the same time it is obvious that $\hat{J}^\mu(0)$ can be free in the point form.

The deep inelastic cross-section is fully defined by the hadronic tensor

$$W^{\mu\nu} = \frac{1}{4\pi} \int e^{\imath qx} \langle P', \chi | \hat{J}^\mu(x) \hat{J}^\nu(0) | P', \chi \rangle d^4x \tag{10}$$

where $|P', \chi\rangle$ is the state of the initial nucleon with the four-momentum P' and the internal wave function χ, and q is the 4-momentum of the virtual photon (for simplicity we shall speak about electromagnetic interactions but the same is valid for weak ones). The state $|P', \chi\rangle$ is the eigenstate of the operator $\hat{P}$ with the eigenvalue P' and the eigenstate of the spin operators $\hat{\mathbf{S}}^2$ and $\hat{S}^z$ which are constructed from $\hat{M}^{\mu\nu}$. In particular, $\hat{P}^2|P', \chi\rangle = m^2|P', \chi\rangle$ where m is the nucleon mass.

The structure of the operator $\hat{P}$ in QCD is rather complicated (see e.g. Refs. [11, 12]) but anyway some of the components of $\hat{P}$ necessarily contain a part which describes the interaction of quarks and gluons at large distances where the QCD running coupling constant α_s is large and perturbation theory does not apply. In view of the last relation, this part is responsible for binding of quarks and gluons in the nucleon. We will call this part the nonperturbative one.

Suppose that the Hamiltonian $\hat{P}^0$ contains the nonperturbative part. Then by analogy with the consideration in Sect. 1, we can show that in the instant form all the operators $\hat{M}^{0i}$ inevitably depend on the nonperturbative part and in the point form all the operators $\hat{P}^k$ inevitably depend on that part. Analogously if the front form Hamiltonian $\hat{P}^-$ contains the nonperturbative part then all the operators $\hat{M}^{-j}$ inevitably depend on this part. Therefore, some components of the operator $\hat{J}^\mu(0)$ in the instant and front forms inevitably depend on the nonperturbative part, and if $\hat{J}^\mu(0) = J^\mu(0)$ in the point form then, as follows from Eq. (3), the operator $\hat{J}^\mu(x)$ in that form inevitably depend on the nonperturbative part. The fact that the same operators $(\hat{P}^\mu, \hat{M}^{\mu\nu})$ describe the transformations of both the operator $\hat{J}^\mu(x)$ and the state $|P', \chi\rangle$ guaranties that $W^{\mu\nu}$ has the correct transformation properties.

We see that the relation between the current operator and the state of the initial nucleon is highly nontrivial. Meanwhile in the present theory they are considered separately. In the framework of the approach to DIS based on Feynman diagrams the possibility of the separate consideration follows from the factorization theorem [13, 14, 15, 16] which asserts in particular that the amplitude of the lepton-parton interaction entering into diagrams dominating in DIS depends only on the hard part of this interaction. Moreover, in leading order in $1/Q$, where $Q = |q^2|^{1/2}$, one obtains the parton model up to anomalous dimensions and perturbative QCD corrections which depend on $\alpha_s(Q^2)$. We will not discuss here the assumptions used in this approach but note that the struck quark is essentially off-shell and its momentum $\tilde{p}$ is usually such that $\tilde{p}^2$ is of order Λ^2 where Λ is the typical hadron mass. Therefore it is far from being obvious that interactions of the struck quark can be considered perturbatively. As noted in Ref. [15], "it is fair to say that a rigorous treatment of factorization has yet to be provided".

For example, in Ref. [14] the proof of the factorization theorem is based on the expansion of the lepton-quark amplitude near the point $\tilde{p}^2 = 0$ (assuming

that the quarks are massless) and only a finite number of terms are taken into account. The result is justified a *posteriori* since it is in agreement with the intuition provided by the parton model. Indeed, in the parton model the quarks are free and therefore $\tilde{p}^2 = 0$. However such a proof by no means excludes a possibility that actually the $\tilde{p}^2$ dependence is important even in leading order in $1/Q$ and therefore the correct result can be obtained only if a sum of an infinite number of diagrams is calculated.

Let us now discuss the following question. Since the current operator depends on the nonperturbative part then this operator depends on the integrals from the quark and gluon field operators over the region of large distances where α_s is large. Is this property compatible with locality? In the framework of canonical formalism compatibility is obvious but, as noted above, the results based on canonical formalism are not reliable. Therefore it is not clear whether in QCD it is possible to construct local electromagnetic and weak current operators beyond perturbation theory. However the usual motivation of the parton model is that, as a consequence of asymptotic freedom (which implies that $\alpha_s(Q^2) \to \infty$ when $Q^2 \to \infty$), the partons in the infinite momentum frame (IMF) are almost free and therefore, at least in leading order in $1/Q$, the nonperturbative part of $\hat{J}^\mu(x)$ is not important.

We will now consider whether this property can be substantiated in the framework of the operator product expansion (OPE) developed by Wilson and others [17]. In this framework the product of the currents entering into Eq. (10) can be written symbolically as

$$\hat{J}(x)\hat{J}(0) = \sum_i C_i(x^2) x_{\mu_1} \cdots x_{\mu_n} \hat{O}_i^{\mu_1 \cdots \mu_n} \tag{11}$$

where $C_i(x^2)$ are the c-number Wilson coefficients while the operators $\hat{O}_i^{\mu_1 \cdots \mu_n}$ depend only on field operators and their covariant derivatives at the origin of Minkowski space and have the same form as in perturbation theory. For example, the basis for twist two operators contains in particular

$$\hat{O}_V^\mu = \mathcal{N}\{\hat{\bar{\psi}}(0)\gamma^\mu\hat{\psi}(0)\} \quad \hat{O}_A^\mu = \mathcal{N}\{\hat{\bar{\psi}}(0)\gamma^\mu\gamma^5\hat{\psi}(0)\} \tag{12}$$

It is important to note that the OPE has been proved only in the framework of perturbation theory and its validity beyond that theory is problematic (see the discussion in Ref. [18] and references therein). Therefore if we use Eq. (11) in DIS we have to assume that either nonperturbative effects

are not important to some orders in $1/Q$ and then we can use Eq. (11) only to these orders (see e.g. Ref. [19]) or it is possible to use Eq. (11) beyond perturbation theory. The question also arises whether Eq. (11) is valid in all the forms of dynamics (as it should be if it is an exact operator equality) or only in some forms.

In the point form all the components of $\hat{P}$ depend on the nonperturbative part and therefore, in view of Eq. (3), it is not clear why there is no nonperturbative part in the x dependence of the right hand side of Eq. (11), or if (for some reasons) it is possible to include the nonperturbative part only into the operators $\hat{O}_i$ then why they have the same form as in perturbation theory.

One might think that in the front form the $C_i(x^2)$ will be the same as in perturbation theory due to the following reasons. The value of q^- in DIS is very large and therefore only a small vicinity of the light cone $x^+ = 0$ contributes to the integral (10). The only dynamical component of $\hat{P}$ is $\hat{P}^-$ which enters into Eq. (11) only in the combination $\hat{P}^- x^+$. Therefore the dependence of $\hat{P}^-$ on the nonperturbative part of the quark-gluon interaction is of no importance. These considerations are not convincing since the integrand is a singular function and the operator $\hat{J}^\mu(0)$ in the front form depends on the nonperturbative part, but nevertheless we assume that Eq. (11) in the front form is valid.

If we assume as usual that there is no problem with the convergence of the OPE series then experiment makes it possible to measure each matrix element $\langle P', \chi|\hat{O}_i^{\mu_1\cdots\mu_n}|P', \chi\rangle$. Let us consider, for example, the matrix element $\langle P', \chi|\hat{O}_V^{\mu}|P', \chi\rangle$. It transforms as a four-vector if the Lorentz transformations of $\hat{O}_V^{\mu}$ are described by the operators $\hat{M}^{\mu\nu}$ describing the transformations of $|P', \chi\rangle$, or in other words, by analogy with Eq. (6)

$$[\hat{M}^{\mu\nu}, \hat{O}_V^{\rho}] = -\imath(\eta^{\mu\rho}\hat{O}_V^{\nu} - \eta^{\nu\rho}\hat{O}_V^{\mu}) \tag{13}$$

It is also clear that Eq. (13) follows from Eqs. (3), (4), (10) and (11). Since the $\hat{M}^{-j}$ in the front form depend on the nonperturbative part, the above considerations make it possible to conclude that at least some components $\hat{O}_V^{\mu}$, and analogously some components $\hat{O}_i^{\mu_1\cdot\mu_n}$, also depend on the nonperturbative part. Since Eq. (13) does not contain any x or q dependence, this conclusion has nothing to do with asymptotic freedom and is valid even in leading order in $1/Q$ (in contrast with the statement of the factorization theorem [13, 14, 15, 16]).

Since the operators $\hat{O}_i^{\mu_1 \cdots \mu_n}$ depend on the nonperturbative part, then by analogy with the above considerations we conclude that the operators in Eq. (12) are ill-defined and the correct expressions for them involve integrals from the field operators over large distances where the QCD coupling constant is large. Therefore it is not clear whether the operators $\hat{O}_i^{\mu_1 \cdots \mu_n}$ are local and whether the Taylor expansion at $x = 0$ is correct, but even it is, the expressions for $\hat{O}_i^{\mu_1 \cdots \mu_n}$ will depend on higher twist operators which contribute even in leading order in $1/Q$.

To understand whether the OPE is valid beyond perturbation theory several authors (see e.g. Ref. [18] and references therein) investigated some two-dimensional models and came to different conclusions. We will not discuss the arguments of these authors but note that the Lie algebra of the Poincare group for 1+1 space-time is much simpler than for 3+1 one. In particular, the Lorentz group is one-dimensional and in the front form the operator M^{+-} is free. Therefore Eqs. (6) and (13) in the "1+1 front form" do not make it possible to conclude that the operators $\hat{J}^\mu(0)$ and $\hat{O}_V^\mu$ necessarily depend on the nonperturbative part. At the same time the full space reflection P in the 1+1 front form is an independent dynamical transformation, in contrast with the situation in the 3+1 front form (see Sect. 1).

Let us now summarize the results of this section. As follows from the commutation relations between the current operator and the generators of the Poincare group representation, the current operator nontrivially depends on the nonperturbative part of the interaction responsible for binding of quarks and gluons in the nucleon. Then the problem arises whether it is possible to construct a local current operator $\hat{J}^\mu(x)$ beyond perturbation theory and whether the nonperturbative part of the interaction entering into $\hat{J}^\mu(x)$ contributes to DIS in leading order in $1/Q$. Our consideration shows that the dependence of $\hat{J}^\mu(x)$ on the nonperturbative part makes the OPE problematic. Nevertheless we assume that Eq. (11) is valid beyond perturbation theory but no form of the operators $\hat{O}_i^{\mu_1 \cdots \mu_n}$ is prescribed. Then we come to conclusion that the nonperturbative part indeed contributes to DIS even in leading order in $1/Q$.

Since the role of nonperturbative effects cannot be neglected even at very large Q, the calculation of their contribution to DIS cannot be carried out in the framework of QCD (at the present stage of this theory). For this reason the remainder of this paper is devoted to the investigation of DIS in models where all the operators in question can be explicitly constructed and it is

possible to explicitly verify that they satisfy Eqs. (1-5).

3 Representations of the Poincare group for systems of interacting particles

As shown by several authors (see e.g. Refs. [20, 21, 22, 23, 24, 25]), the representation generators satisfying Eqs. (1) and (2) can be explicitly constructed in the framework of relativistic quantum mechanics (RQM) of systems with a fixed number of interacting particles and, as shown in Refs. [26, 27], in this framework it is also possible to explicitly construct the current operator satisfying Eqs. (3-5).

Of course, as already noted, any model consideration cannot have any pretensions to be fundamental. Nevertheless in this paper we argue that the consideration in the framework of RQM can shed light on the problem why the effect of binding in DIS is important even at very large values of Q.

For the consideration of a system of interacting particles it is useful to consider first kinematical relations in the corresponding system of free particles and then consider the introduction of the interaction into this system.

Consider a system of N free particles with the 4-momenta p_i, masses m_i, spins s_i, and the electric charges e_i $(i = 1, ...N)$. We also use σ_i to denote the projection of the spin of particle i on the z axis and use $\perp$ to denote the projection of the three-dimensional vectors onto the plane xy. For definiteness we assume that $m_i > 0$, but this assumption is not crucial. Formally the subsequent consideration will not depend on whether N is finite or infinite but in the last case essential efforts should be made to prove that all the operators in question are well defined. For this reason we assume that N is final.

The Hilbert space H for the system under consideration is the space of functions $\varphi(\mathbf{p}_{1\perp}, p_1^+, \sigma_1, ...\mathbf{p}_{N\perp}, p_N^+, \sigma_N)$ such that

$$\sum_{\sigma_1...\sigma_N} \int |\varphi(\mathbf{p}_{1\perp}, p_1^+, \sigma_1, ...\mathbf{p}_{N\perp}, p_N^+, \sigma_N)|^2 \prod_{i=1}^{N} d\rho(\mathbf{p}_{i\perp}, p_i^+) \quad < \infty \tag{14}$$

where

$$d\rho(\mathbf{p}_\perp, p^+) = \frac{d^2\mathbf{p}_\perp dp^+}{2(2\pi)^3 p^+} \tag{15}$$

We define $P = p_1 + ... + p_N$, $M_0 = |P| \equiv (P^2)^{1/2}$, and $G = PM_0^{-1}$. Let $\beta(G) \equiv \beta(\mathbf{G}_\perp, G^+) \in SL(2,C)$ be the matrix with the components

$$\begin{aligned}
&\beta_{11} = \beta_{22}^{-1} = 2^{1/4}(G^+)^{1/2}, \quad \beta_{12} = 0, \\
&\beta_{21} = (G^x + \imath G^y)\beta_{22},
\end{aligned} \tag{16}$$

$L(\beta)$ be the Lorentz transformation corresponding to $\beta \in SL(2,C)$ and

$$k_i = L[\beta(G)]^{-1}p_i \quad (i = 1, ...N) \tag{17}$$

The four-vectors p_i have the components $(\omega_i(\mathbf{p}_i), \mathbf{p}_i)$, and the four-vectors k_i have the components $(\omega_i(\mathbf{k}_i), \mathbf{k}_i)$ where $\omega_i(\mathbf{k}) = (m_i^2 + \mathbf{k}^2)^{1/2}$. In turn, only $N-1$ vectors $\mathbf{k}_i$ are independent since, as follows from Eqs. (16) and (17), $\mathbf{k}_1 + ... + \mathbf{k}_N = 0$. Therefore $L[\beta(G)]$ has the meaning of the boost, and $\mathbf{k}_i$ are the momenta in the c.m. frame. It is easy to show that $M_0 = \omega_1(\mathbf{k}_1) + ... + \omega_N(\mathbf{k}_N)$.

Let us define

$$d\rho^F(int) = 2(2\pi)^3 M_0 \delta^{(3)}(\mathbf{k}_1 + \cdots + \mathbf{k}_N) \prod_{i=1}^{N} d\rho(\mathbf{k}_{i\perp}, k_i^+) \tag{18}$$

We also define the "internal' space H_{int}^F as the space of functions $\chi^F(\mathbf{k}_1, \sigma_1, ...\mathbf{k}_N, \sigma_N)$ such that

$$||\chi^F||_F^2 = \sum_{\sigma_1...\sigma_N} \int |\chi(\mathbf{k}_1, \sigma_1, ...\mathbf{k}_N, \sigma_N)|^2 d\rho^F(int) \quad < \infty \tag{19}$$

Then the space of functions satisfying Eq. (14) can be realized as the space of functions $\varphi(\mathbf{P}_\perp, P^+)$ with the range in H_{int}^F and such that

$$\int ||\varphi(\mathbf{P}_\perp, P^+)||_F^2 d\rho(\mathbf{P}_\perp, P^+) \quad < \infty \tag{20}$$

For noninteracting particles the representation generators of the Poincare group are equal to sums of the corresponding one-particle generators. A direct calculation of these sums shows that in the variables $\mathbf{P}_\perp, P^+, \mathbf{k}_1, ...\mathbf{k}_N$ the generators have the form

$$P^+ = P^+, \quad \mathbf{P}_\perp = \mathbf{P}_\perp, \quad P^- = \frac{M_0^2 + \mathbf{P}_\perp^2}{2P^+},$$

$$M^{+-}=\imath P^+\frac{\partial}{\partial P^+},\quad M^{+j}=-\imath P^+\frac{\partial}{\partial P^j},$$
$$M^{xy}=l^z(\mathbf{P}_\perp)+S^z,\quad M^{-j}=-\imath(P^j\frac{\partial}{\partial P^+}+P^-\frac{\partial}{\partial P^j})-$$
$$\frac{\epsilon_{jl}}{P^+}(M_0S^l+P^lS^z) \tag{21}$$

Here the first expression implies that the generator P^+ is equal to the operator of multiplication by the variable P^+ defined above and the second expression should be understood analogously. The indices j,l take the values $1,2$ and ϵ_{jl} is the antisymmetric tensor with the components $\epsilon_{12}=-\epsilon_{21}=1$, $\epsilon_{11}=\epsilon_{22}=0$. We also use $\mathbf{l}(\mathbf{P})=-\imath\mathbf{P}\times(\partial/\partial\mathbf{P})$ to denote the orbital angular-momentum operator and $\mathbf{S}$ to denote the N-particle spin operator. The latter acts only through the variables of the space H_{int}^F and the explicit form of $\mathbf{S}$ is of no importance for us. For $N=2$ and $N=3$ this form and a detailed derivation of Eq. (21) can be found, for example, in Refs. [28, 29, 25]; the generalization to the case of arbitrary N is obvious.

As follows from Eq. (17)

$$\xi_i\equiv\frac{p_i^+}{P^+}=\frac{\omega_i(\mathbf{k}_i)+k_i^z}{M_0(\mathbf{k}_1,...\mathbf{k}_N)}\in(0,1) \tag{22}$$

Therefore $(\mathbf{k}_{1\perp},\xi_1,...\mathbf{k}_{N\perp},\xi_N)$ also is a possible choice of the internal momentum variables, and these variables are constrained by the relations $\mathbf{k}_{1\perp}+...+\mathbf{k}_{N\perp}=0$, $\xi_1+...+\xi_N=1$.

As follows from Eq. (18), if $N=2$ then

$$d\rho^F(int)=\frac{d^2\mathbf{k}_\perp d\xi}{2(2\pi)^3\xi(1-\xi)}=\frac{M_0(\mathbf{k})d^3\mathbf{k}}{2(2\pi)^3\omega_1(\mathbf{k})\omega_2(\mathbf{k})} \tag{23}$$

where $\mathbf{k}\equiv\mathbf{k}_1=-\mathbf{k}_2$, $\xi=\xi_1=1-\xi_2$. Let int_i be a full set of the internal momentum variables for the system $(1,...i-1,i+1,...N)$ and $d\rho^F(int_i)$ be the internal volume element in the internal space $H_{int}^{(i)F}$ for this system. Then, by analogy with the derivation of Eq. (23) from Eq. (18), it is easy to show that

$$d\rho^F(int)=\frac{d^2\mathbf{k}_{i\perp}d\xi_i}{2(2\pi)^3\xi_i(1-\xi_i)}d\rho^F(int_i) \tag{24}$$

Therefore the normalization condition (19) can be written as

$$||\chi^F||_F^2=\sum_{\sigma_1...\sigma_N}\int|\chi^F(\mathbf{k}_{i\perp},\xi_i,int_i,\sigma_1,...\sigma_N)|^2\cdot$$

$$\frac{d^2\mathbf{k}_{i\perp} d\xi_i}{2(2\pi)^3\xi_i(1-\xi_i)} d\rho^F(int_i) \, < \, \infty \tag{25}$$

If M_i is the free mass operator for the system $(1, ...i-1, i+1, ...N)$ then, as follows from Eqs. (17) and (22),

$$M_0 \equiv M_0(\mathbf{k}_i, M_i) \equiv M_0(\mathbf{k}_{i\perp}, \xi_i, M_i) = \omega_i(\mathbf{k}_i) + \\ (M_i^2 + \mathbf{k}_i^2)^{1/2} = [\frac{m_i^2}{\xi_i} + \frac{M_i^2}{1-\xi_i} + \frac{\mathbf{k}_{i\perp}^2}{\xi_i(1-\xi_i)}]^{1/2} \tag{26}$$

Instead of $(\mathbf{P}_\perp, P^+)$ it is also possible to choose $(\mathbf{G}_\perp, G^+)$ as the external variables while the internal variables can be chosen as above. Now we introduce H_{int}^P as the space of functions $\chi^P(\mathbf{k}_1, \sigma_1, ...\mathbf{k}_N, \sigma_N)$ such that

$$||\chi^P||_P^2 = \sum_{\sigma_1 ... \sigma_N} \int |\chi^P(\mathbf{k}_1, \sigma_1, ...\mathbf{k}_N, \sigma_N)|^2 d\rho^P(int) \, < \, \infty \tag{27}$$

where

$$d\rho^P(int) = M_0^2 d\rho^F(int) \tag{28}$$

Then by analogy with Eq. (20) it is easy to show that the Hilbert space of functions satisfying Eq. (14) can be realized as the space of functions $\varphi(\mathbf{G}_\perp, G^+)$ with the range in H_{int}^P and such that

$$\int ||\varphi(\mathbf{G}_\perp, G^+)||_P^2 d\rho(\mathbf{G}_\perp, G^+) \, < \, \infty \tag{29}$$

A direct calculation shows that in the variables $(\mathbf{G}_\perp, G^+, \mathbf{k}_1, ...\mathbf{k}_N)$ the representation generators of the Poincare group have the form (compare with Eq. (21))

$$P^+ = M_0 G^+, \quad \mathbf{P}_\perp = M_0 \mathbf{G}_\perp, \quad P^- = M_0 G^- = M_0 \frac{1+\mathbf{G}_\perp^2}{2G^+}, \\ M^{+-} = \imath G^+ \frac{\partial}{\partial G^+}, \quad M^{+j} = -\imath G^+ \frac{\partial}{\partial G^j}, \\ M^{xy} = l^z(\mathbf{G}_\perp) + S^z, \quad M^{-j} = -\imath(G^j \frac{\partial}{\partial G^+} + G^- \frac{\partial}{\partial G^j}) - \\ \frac{\epsilon_{jl}}{G^+}(S^l + G^l S^z) \tag{30}$$

If the particles interact with each other then one of the simplest way to preserve the relativistic commutation relations (1) is to replace M_0 in Eq.

(21) by the mass operator $\hat{M}^F$ which acts only through the variables of the space H_{int}^F and commutes with $\mathbf{S}$. Then the representation generators of the Poincare group obviously have the form

$$
\begin{aligned}
&P^+ = P^+, \quad \mathbf{P}_\perp = \mathbf{P}_\perp, \quad \hat{P}^- = \frac{(\hat{M}^F)^2 + \mathbf{P}_\perp^2}{2P^+}, \\
&M^{+-} = \imath P^+ \frac{\partial}{\partial P^+}, \quad M^{+j} = -\imath P^+ \frac{\partial}{\partial P^j}, \\
&M^{xy} = l^z(\mathbf{P}_\perp) + S^z, \quad \hat{M}^{-j} = -\imath(P^j \frac{\partial}{\partial P^+} + \hat{P}^- \frac{\partial}{\partial P^j}) - \\
&\frac{\epsilon_{jl}}{P^+}(\hat{M}^F S^l + P^l S^z)
\end{aligned}
\qquad (31)
$$

Such a procedure was first proposed by Bakamdjian and Thomas [30]. It is obvious that the generators in Eq. (31) are given in the front form.

Analogously, we can replace M_0 in Eq. (30) by the mass operator $\hat{M}^P$ which acts only through the variables of the space H_{int}^P and commutes with $\mathbf{S}$. Then the representation generators of the Poincare group obviously have the form

$$
\begin{aligned}
&\hat{P}^+ = \hat{M}^P G^+, \quad \hat{\mathbf{P}}_\perp = \hat{M}^P \mathbf{G}_\perp, \quad \hat{P}^- = \hat{M}^P G^- = \\
&\hat{M}^P \frac{1 + \mathbf{G}_\perp^2}{2G^+}, \quad M^{+-} = \imath G^+ \frac{\partial}{\partial G^+}, \quad M^{+j} = -\imath G^+ \frac{\partial}{\partial G^j}, \\
&M^{xy} = l^z(\mathbf{G}_\perp) + S^z, \quad M^{-j} = -\imath(G^j \frac{\partial}{\partial G^+} + G^- \frac{\partial}{\partial G^j}) - \\
&\frac{\epsilon_{jl}}{G^+}(S^l + G^l S^z)
\end{aligned}
\qquad (32)
$$

The generators in this expression are obviously given in the point form.

If Γ_i^F and Γ_i^P $(i = 1, ...10)$ are the generators given by Eqs. (31) and (32) respectively, then, as shown by Sokolov and Shatny [31], if $\hat{M}^F$ and $\hat{M}^P$ are unitarily equivalent, the sets Γ_i^F and Γ_i^P are unitarily equivalent too.

In addition to extended relativistic invariance, the generators should also satisfy cluster separability (see Refs. [20, 21, 22] for details). In the framework of the method of packing operator developed by Sokolov [20] it can be shown [20, 21, 23] that the most general form of the generators in the front and point forms is $A^F \Gamma_i^F (A^F)^{-1}$ and $A^P \Gamma_i^P (A^P)^{-1}$ respectively where A^F and A^P are some unitary operators.

We shall consider a simple model when $\hat{M}^P = M_0 + v_N^P$ where v_N^P is the fully linked part of the mass operator. In other words, there is only the N-body interaction and there are no N'-body interactions if $N' < N$. Then relativistic invariance and cluster separability are obviously satisfied if $[v_N^P, \mathbf{S}] = 0$ and $A^P = 1$. Analogously we can consider the case when $\hat{M}^F = M_0 + v_N^F$, $[v_N^F, \mathbf{S}] = 0$ and $A^F = 1$.

Of course, the real mass operator contains all N'-body interactions and $A^P \neq 1$. However, it is reasonable to assume that the nucleon binding energy and wave function are mainly defined by some confining interaction v_N which cannot be determined in the framework of perturbative QCD, while the other interaction operators in $\hat{M}^P$ can be determined in this framework. Such a situation takes place in many realistic models which successfully describe the baryon spectra (see e.g. Ref. [32]).

4 Parton model as a consequence of IA in the front form of dynamics

For systems with a finite number of particles the operator $\hat{J}^\mu(x)$ as a usual operator valued function of x is well defined [27] and, as follows from Eqs. (1), (3) and (4), if the operator $\hat{P}$ is known then $\hat{J}^\mu(x)$ is fully defined by the operator $\hat{J}^\mu(0)$ satisfying Eq. (6).

We shall always assume that all particles having the electric charge are structureless and their spin is equal to 1/2. Then the one-particle current operator for particle i acts over the variables of this particle as

$$\begin{aligned} J_i^\mu(0)\varphi(\mathbf{p}_{i\perp}, p_i^+, \sigma_i) \quad = \quad & r_i \sum_{\sigma_i'} \int [\bar{w}_i(p_i, \sigma_i)\gamma^\mu w_i(p_i', \sigma_i')] \cdot \\ & \varphi(\mathbf{p}_{i\perp}', p_i^{'+}, \sigma_i') d\rho(\mathbf{p}_{i\perp}', p_i^{'+}) \end{aligned} \tag{33}$$

and over the variables of other particles it acts as the identity operator. Here $r_i = e_i/e_0$ is the ratio of the particle electric charge to the unit electric charge, $w_i(p_i, \sigma_i)$ is the Dirac light cone spinor, γ^μ is the Dirac γ-matrix, and $\bar{w} = w^+\gamma^0$. The form of $w_i(p_i, \sigma_i)$ in the spinor representation of the Dirac γ-matrices is

$$w_i(p_i, \sigma_i) = \sqrt{m_i} \left\| \begin{array}{c} \beta(\mathbf{p}_{i\perp}/m_i, p_i^+/m_i)\chi(\sigma_i) \\ \beta(\mathbf{p}_{i\perp}/m_i, p_i^+/m_i)^{-1+}\chi(\sigma_i) \end{array} \right\| \tag{34}$$

where $\chi(\sigma)$ is the ordinary spinor describing the state with the spin projection on the z axis equal to σ.

By definition, the operator $\hat{J}^\mu(0)$ in IA is given by

$$\hat{J}^\mu(0) = J^\mu(0) = \sum_{i=1}^{N} J_i^\mu(0) \tag{35}$$

In the front form such a current operator satisfies neither Lorentz invariance nor invariance under the discrete symmetries. The violation of Lorentz invariance has been explained in Sect. 2 while the violation of the discrete symmetries follows from the relations

$$\begin{aligned}&\hat{U}_P(\hat{J}^0(0), \hat{\mathbf{J}}(0))\hat{U}_P^{-1} = (\hat{J}^0(0), -\hat{\mathbf{J}}(0)) = \\ &\hat{U}_T(\hat{J}^0(0), \hat{\mathbf{J}}(0))\hat{U}_T^{-1}\end{aligned} \tag{36}$$

(recall that $\hat{U}_P$ and $\hat{U}_T$ in the front form are necessarily interaction dependent). Nevertheless in this section we carry out the calculations in IA for $\hat{J}^\mu(0)$ in the front form and show that the results are exactly the same as in the parton model. The analogous calculations have been carried out elsewhere (see e.g. Ref. [33, 34] and references therein).

As follows from Eqs. (3) and (10), the hadronic tensor can be written in the form

$$\begin{aligned}W^{\mu\nu} \;=\;& \frac{1}{4\pi}\sum_X (2\pi)^4 \delta^{(4)}(P' + q - P)\langle P', \chi^F| \\ & J^\mu(0)|X\rangle\langle X|J^\nu(0)|P', \chi^F\rangle\end{aligned} \tag{37}$$

where the sum is taken over all possible final states $|X\rangle$, and P is the four-momentum of the state $|X\rangle$. We will calculate this tensor only in the Bjorken limit when Q^2 and $P'q$ are very large but the Bjorken variable $x = Q^2/2(P'q)$ is not too close to 0 or 1 (from now on we will use x only to denote the Bjorken variable).

We assume as usual (see e.g. Refs. [13, 14, 35, 36]) that the final state interaction of the struck quark with the remnants of the target is the effect of order $(m/Q)^2$. Then in the model for the mass operator considered in Sect. 3, we can write $|X\rangle$ as the states of N free particles with the 4-momenta p_i" and spin projections σ_i" $(i = 1, ...N)$:

$$|X\rangle = \prod_{i=1}^{N} |p_i", \sigma_i"\rangle \tag{38}$$

Taking into account the normalization of free states in the scattering theory we write the wave functions of these states in the form

$$|p_i", \sigma_i"\rangle = 2(2\pi)^3 p_i^{"+} \delta^{(2)}(\mathbf{p}_{i\perp} - \mathbf{p}_{i\perp}")\delta(p_i^+ - p_i^{"+})\delta_{\sigma_i \sigma_i"} \tag{39}$$

where $\delta_{\sigma_i \sigma_i"}$ is the Cronecker symbol.

Analogously if the Poincare group generators are given by Eq. (31), the wave function of the initial nucleon can be written as

$$|P', \chi^F\rangle = 2(2\pi)^3 P'^+ \delta^{(2)}(\mathbf{P}_\perp - \mathbf{P}'_\perp)\delta(P^+ - P'^+)\chi^F \tag{40}$$

where the internal nucleon wave function χ^F is normalized as $||\chi^F||_F = 1$.

From now on it will be convenient to denote the momenta of final particles without two primes, i.e. as $p_1, ...p_N$. Then, as follows from Eqs. (33), (35) and (37-40),

$$\begin{aligned}
&W^{\mu\nu} = \frac{1}{4\pi} \sum_{\sigma_1 ... \sigma_N} \sum_{i,j=1}^{N} \sum_{\sigma'_i, \sigma'_j} \int \{\prod_{l=1}^{N} d\rho(\mathbf{p}_{l\perp}, p_l^+)\} \frac{r_i r_j}{\xi'_i \xi'_j} \cdot \\
&(2\pi)^4 \delta^{(4)}(P' + q - P)\langle \chi^F(\mathbf{k}'_j, int_j, \sigma_1, ...\sigma'_j, ...\sigma_N)| \\
&[\bar{w}_j(p'_j, \sigma'_j)\gamma^\mu w_j(p_j, \sigma_j)][\bar{w}_i(p_i, \sigma_i)\gamma^\nu w_i(p'_i, \sigma'_i)] \\
&|\chi^F(\mathbf{k}'_i, int_i, \sigma_1, ...\sigma'_i, ...\sigma_N)\rangle
\end{aligned} \tag{41}$$

where p'_i, $\mathbf{k}'_i$ and ξ'_i are defined as follows.

If $P_i = p_1 + ... + p_{i-1} + p_{i+1} + ...p_N$ then the four-vector p'_i is such that $p_i'^2 = m_i^2$, $p_i'^+ = P'^+ - P_i^+$, $\mathbf{p}'_{i\perp} = \mathbf{P}'_\perp - \mathbf{P}_{i\perp}$. Then (compare with Eqs. (17) and (22))

$$\begin{aligned}
k'_i \;\; &= \;\; L[\beta(\frac{\mathbf{P}'_\perp}{M_0(\mathbf{k}'_i, M_i)}, \frac{P'^+}{M_0(\mathbf{k}'_i, M_i)})]^{-1} p'_i, \\
&\xi'_i = \frac{\omega_i(\mathbf{k}'_i) + k_i'^z}{M_0(\mathbf{k}'_i, M_i)}
\end{aligned} \tag{42}$$

and $k'_i = (\omega_i(\mathbf{k}'_i), \mathbf{k}'_i)$.

We can write P_i on the one hand as a function of P', p'_i, int_i and on the other hand as a function of P, p_i, int_i. Therefore (compare with Eq. (17)) $\mathbf{k}'_i$

can be defined by the condition

$$L[\beta(\frac{\mathbf{P}'_\perp}{M_0(\mathbf{k}'_i,M_i)},\frac{P'^+}{M_0(\mathbf{k}'_i,M_i)})]((M_i^2+\mathbf{k}_i'^2)^{1/2},-\mathbf{k}'_i) =$$
$$L[\beta(\frac{\mathbf{P}_\perp}{M_0(\mathbf{k}_i,M_i)},\frac{P^+}{M_0(\mathbf{k}_i,M_i)})]((M_i^2+\mathbf{k}_i^2)^{1/2},-\mathbf{k}_i) \tag{43}$$

As follows from Eq. (41), the mass $M = M_0(\mathbf{k}_i, M_i)$ of the final state satisfies the condition $M^2 = (P' + q)^2$. Hence

$$M^2 = m^2 + \frac{Q^2(1-x)}{x} \tag{44}$$

and in the Bjorken limit m^2 in this expression can be neglected.

It is convenient to consider the process in the IMF where $\mathbf{P}'_\perp = \mathbf{q}_\perp = 0$, and P'^z is positive and very large. By analogy with the Breit frame for elastic processes we choose the reference frame in which $\mathbf{P} + \mathbf{P}' = 0$. It is easy to show that in this reference frame

$$q^0 = 2|\mathbf{P}'|(1-x), \quad P'^+ = \sqrt{2}|\mathbf{P}'|, \quad q^+ = -\sqrt{2}|\mathbf{P}'|x,$$
$$P^+ = \sqrt{2}|\mathbf{P}'|(1-x) \tag{45}$$

Then as follows from Eqs. (16), (22), (43) and (45)

$$\mathbf{k}'_{i\perp} = \mathbf{k}_{i\perp}, \quad \xi'_i = (1-x)\xi_i + x \tag{46}$$

We assume that the internal wave function χ^F effectively cuts the contribution of large momenta, and therefore the contribution to the integrals containing $\chi^F(\mathbf{k}'_i, int_i)$ is given only by the momenta with $|\mathbf{k}'_i| \leq m_0$, $|\mathbf{k}_l| \leq m_0$ $(l = 1, ...i-1, i+1, ...N)$ where m_0 is some parameter satisfying the condition $m_0 \ll Q$. In turn, as follows from Eq. (42), this implies that ξ'_i should not be close to 0 or 1, and, as follows from Eq. (26), in the Bjorken limit $|k_i^z| \approx M/2$. If $k^z > 0$ then, as follows from Eqs. (22) and (26), $\xi_i \approx 1$. However in this case ξ'_i is close to 1 as follows from Eq. (46). Therefore the only possibility is $k^z < 0$, $\xi_i \approx 0$. Then, as follows from Eq. (46), in the Bjorken limit $\xi'_i = x$. Analogously in the Bjorken limit the contribution to the integrals containing $\chi^F(\mathbf{k}'_j, int_j)$ is not negligible only if $|\mathbf{k}'_j| \leq m_0$, $|\mathbf{k}_n| \leq m_0$ $(n = 1, ...j-1, j+1, ...N)$, $|k_j^z| \approx M/2$ and $k_j^z < 0$. All these conditions are

compatible with each other only if $i = j$, i.e. the quarks absorb the virtual photon incoherently. As follows from Eqs. (22) and (42), the result $\xi_i' = x$ fully agrees with the interpretation of the quantity x in the parton model as the momentum fraction of the struck quark in the IMF.

As follows from Eqs. (20), (24) and (26), if the above conditions are satisfied then

$$(2\pi)^4\delta^{(4)}(P'+q-P)\prod_{i=1}^{N} d\rho(\mathbf{p}_{i\perp},p_i^+) = \frac{d^2\mathbf{k}_{i\perp}}{8\pi^2 M^2} d\rho^F(int_i) \qquad (47)$$

Therefore, as follows from Eqs. (17), (34), (41), and (45-47), the final result for the hadronic tensor is

$$\begin{aligned} W^{\mu\nu} \;&=\; \sum_{i=1}^{N} r_i^2 \int \langle \chi^F(\mathbf{k}_{i\perp},\xi_i'=x,int_i)|S_i^{\mu\nu}| \\ &\quad \chi^F(\mathbf{k}_{i\perp},\xi_i'=x,int_i)\rangle \frac{d^2\mathbf{k}_{i\perp} d\rho^F(int_i)}{4(2\pi)^3 x(1-x)} \end{aligned} \qquad (48)$$

where we do not write the spin variables in the arguments of the function χ^F, the scalar product is taken over these variables, and the tensor operator $S_i^{\mu\nu}$ is as follows. It is equal to zero if either μ or ν is equal to $\pm$, while if $j,l = x,y$ then $S_i^{jl} = \delta_{jl} + 2\imath\epsilon_{jl}s_i^z$, where s_i^z is the z component of the spin operator for particle i.

Let us introduce the notation

$$\rho_i(x) = \sum_{\sigma_1...\sigma_N} \int |\chi^F(\mathbf{k}_{i\perp},\xi_i'=x,int_i,\sigma_1,...\sigma_N)|^2 \frac{d^2\mathbf{k}_{i\perp} d\rho^F(int_i)}{2(2\pi)^3 x(1-x)} \qquad (49)$$

This quantity is called the parton density since, as follows from Eq. (25), $\rho_i(\xi_i')d\xi_i'$ is the probability of the event that particle i in the bound state has the value of ξ_i' in the interval $(\xi_i',\xi_i'+d\xi_i')$.

It is well-known that the average value of the hadronic tensor over all initial spin states is equal to

$$\begin{aligned} &W^{\mu\nu}(P',q) = (\frac{q^\mu q^\nu}{q^2} - \eta^{\mu\nu})F_1(x,q^2) + \\ &\frac{1}{(P'q)}(P'^\mu - \frac{q^\mu(P'q)}{q^2})(P'^\nu - \frac{q^\nu(P'q)}{q^2})F_2(x,q^2), \end{aligned} \qquad (50)$$

Then, as follows from Eqs. (48) and (49), the structure functions F_1 and F_2 depend only on x (this phenomenon is known as Bjorken Scaling):

$$F_1(x) = \frac{1}{2}\sum_{i=1}^{N} r_i^2 \rho_i(x), \quad F_2(x) = 2xF_1(x) \tag{51}$$

(the last equality is known as the Callan-Gross relation [37]). These expressions for the structure functions have been derived by many authors in the framework of the parton model. Equation (48) also makes it possible to write the expression for the polarized structure functions, but we shall not dwell on this question.

One might think that the above results are natural since they fully agree with the parton model. However the following question arises. Since the current in IA does not satisfy Lorentz invariance, P invariance and T invariance, the results for the structure functions depend on the reference frame in which these functions are calculated. An argument in favor of choosing the IMF is that in this reference frame the current conservation in the Bjorken limit is restored since the tensor $W^{\mu\nu}$ given by Eq. (48) satisfies the continuity equation $q_\mu W^{\mu\nu} = q_\nu W^{\mu\nu} = 0$. Another well-known arguments are based on the approach proposed by Weinberg [38] and developed by several authors (see, for example, Refs. [39, 40]). Let us note however that though quantum field theory in the IMF seems natural and has some advantages, it also has some serious difficulties which are not present in the usual formulation [41].

In our opinion, a rather strange feature of the above results is as follows. By looking through the derivation of these results one can easily see that the initial state is treated in fact not as the bound state but as the free state of noninteracting particles. Indeed, we have never used the fact that the initial state is the eigenstate of the mass operator $\hat{M}^F$ with the eigenvalue m: $\hat{M}^F\chi^F = m\chi^F$. In IA in the front form the relation between the quantities $\mathbf{k}'_i$ and $\mathbf{k}_i$ (see Eq. (43)) is derived from the condition that the four-vectors $(M_i^2 + \mathbf{k}_i'^2)^{1/2}, -\mathbf{k}'_i)$ and $(M_i^2 + \mathbf{k}_i^2)^{1/2}, -\mathbf{k}_i)$ are connected by the Lorentz boosts in the initial and final states. The problem whether particle i does not interact with the other particles in the final state deserves a separate investigation and we will not discuss this problem in the present paper, but it is strange that we neglect the interaction in the initial state and write the free mass $M_0(\mathbf{k}'_i, M_i)$ instead of the real mass m which has the initial state.

The effect of binding can be explicitly taken into account in models where

the current operator satisfies relativistic invariance and current conservation. This problem is considered in the next section.

The unpolarized hadronic tensor for the neutrino (antineutrino) - nucleon scattering has the form

$$
\begin{aligned}
W^{\mu\nu}(P',q) \;&=\; -\eta^{\mu\nu}F_1(x,q^2)+\frac{P'^{\mu}P'^{\nu}}{2(P'q)}F_2(x,q^2)- \\
&\qquad \imath\epsilon^{\mu\nu\rho\lambda}\frac{P'_{\rho}q_{\lambda}}{2(P'q)}F_3(x,q^2)
\end{aligned}
\tag{52}
$$

where $\epsilon^{\mu\nu\rho\lambda}$ is the fully antisymmetric tensor with $\epsilon^{0123}=1$ and the terms with q^{μ} and q^{ν} are dropped since they give zero after multiplication by the leptonic tensor. The calculation of the deep inelastic scattering caused by the charged weak currents can be carried out by analogy with the above calculation and the result is

$$
\begin{aligned}
&F_1^{\nu p}(x)=\sum_{i=d,s,\bar{u},\bar{c}}\rho_i(x), \quad F_1^{\bar{\nu}p}(x)=\sum_{i=u,c,\bar{d},\bar{s}}\rho_i(x),\\
&F_3^{\nu p}(x)=2\sum_{i=d,s}\rho_i(x)-2\sum_{i=\bar{u},\bar{c}}\rho_i(x),\\
&F_3^{\bar{\nu}p}(x)=2\sum_{i=u,c}\rho_i(x)-2\sum_{i=\bar{d},\bar{s}}\rho_i(x),\\
&F_2^{\nu p}(x)=2xF_1^{\nu p}(x), \quad F_2^{\bar{\nu}p}(x)=2xF_1^{\bar{\nu}p}(x)
\end{aligned}
\tag{53}
$$

where we assume for simplicity that the proton does not contain b and t quarks.

5 Consistent calculation of the hadronic tensor

In this section we calculate the hadronic tensor assuming that $\hat{J}^{\mu}(0)=J^{\mu}(0)$ in the point form. Then extended Poincare invariance is satisfied automatically (see Sects. 1 and 2) and if $\hat{J}^{\mu}(0)$ is free in some reference frame it will remain free in all reference frames obtained from this one by means of Lorentz boosts. As shown in Ref. [27], the operator $\hat{J}^{\mu}(0)$ is fully defined by its matrix elements in the reference frame where $\mathbf{G}'+\mathbf{G}=0$. Here $G'=P'/m$

is the four-velocity of the system in the initial state and $G = P/M$ is the same quantity in the final state.

Since now the current operator satisfies Poincare invariance the calculations can be carried out in any reference frame. Therefore we can again suppose that $\mathbf{P}'_\perp = \mathbf{q}_\perp = 0$ and $P'^z > 0$. Therefore $\mathbf{G}'_\perp = \mathbf{G}_\perp = 0$, and $G'^z > 0$. If $\mathbf{G} + \mathbf{G}' = 0$ then $G^z < 0$. As follows from these expressions and Eq. (44), in the Bjorken limit

$$(G^z)^2 = \frac{Q}{4m[x(1-x)]^{1/2}} \tag{54}$$

Therefore $G'^z \gg 1$ in the Bjorken limit.

As follows from Eq. (5), in the reference frame under consideration the $\perp$ components of $\hat{J}^\mu(0)$ are not constrained by the continuity equation while the longitudinal components should satisfy the condition

$$\begin{aligned}&(M-m)G'^0\langle X|\hat{J}^0(0)|P',\chi^P\rangle = \\ &-(M+m)G'^z\langle X|\hat{J}^z(0)|P',\chi^P\rangle\end{aligned} \tag{55}$$

Therefore in the Bjorken limit

$$\langle X|\hat{J}^0(0)|P',\chi^P\rangle = -\langle X|\hat{J}^z(0)|P',\chi^P\rangle \tag{56}$$

Note that the operator G is free by construction (see Eq. (32)) and therefore the free operator $J^\mu(0)$ satisfies the condition

$$\begin{aligned}&G'^0\langle X|J^0(0)(M-M_0)|P',\chi^P\rangle = \\ &-G'^z\langle X|J^z(0)(M+M_0)|P',\chi^P\rangle\end{aligned} \tag{57}$$

Since $M \gg M_0$ in the Bjorken limit we conclude that the choice $\hat{J}^\mu(0) = J^\mu(0)$ is compatible with the continuity equation.

This choice is also compatible with macrolocality. Indeed, as follows from Eqs. (3) and (32)

$$\begin{aligned}&\hat{J}^\mu(\mathbf{x})\hat{J}^\nu(0)\varphi(G) = exp(-\imath\hat{M}^P(\mathbf{G}\mathbf{x}))J^\mu(0)\cdot \\ &exp(\imath\hat{M}^P(\mathbf{G}\mathbf{x}))J^\nu(0)\varphi(G)\end{aligned} \tag{58}$$

Since the operator $J^\mu(0)$ is not singular, the strong limit of $\hat{J}^\mu(\mathbf{x})\hat{J}^\nu(0)$ is equal to zero if $|\mathbf{x}| \to \infty$. This can be proved, for example, by analogy with

the proof of space separability in Ref. [42]. Analogously it is easy to see that the same is valid for the strong limit of $\hat{J}^\nu(0)\hat{J}^\mu(\mathbf{x})$.

We conclude that the choice $\hat{J}^\mu(0) = J^\mu(0)$ in the point form is consistent in the sense that it is compatible with extended Poincare invariance, continuity equation and macrolocality.

As follows from Eq. (32) and the normalization of states in the scattering theory, the wave function of the initial nucleon in the point form can be written as

$$|P',\chi^P\rangle = \frac{2}{m}(2\pi)^3 G'^+\delta^{(2)}(\mathbf{G}_\perp - \mathbf{G}'_\perp)\delta(G^+ - G'^+)\chi^P \tag{59}$$

where $||\chi^P||_P = 1$.

Now the hadronic tensor has the form (compare with Eq. (37))

$$\begin{aligned} W^{\mu\nu} \;\; = \;\; & \frac{1}{4\pi}\sum_X (2\pi)^4\delta^{(4)}(P'+q-P)\langle P',\chi^P| \\ & J^\mu(0)|X\rangle\langle X|J^\nu(0)|P',\chi^P\rangle \end{aligned} \tag{60}$$

We will calculate this tensor in the reference frame considered above. As follows from Eq. (56), in this reference frame the matrix elements of the operator $\hat{J}^+(0)$ is negligible in comparison with the matrix elements of the operator $\hat{J}^-(0)$. Therefore it is sufficient to calculate the tensor $W^{\mu\nu}$ for $\mu,\nu = x,y,-$. It is also easy to show that in this reference frame

$$\mathbf{q}_\perp = 0, \quad q^+ = -\frac{x^{3/4}(mQ)^{1/2}}{\sqrt{2}(1-x)^{1/4}}, \quad q^- = \frac{(1-x)^{1/4}Q^{3/2}}{\sqrt{2}m^{1/2}x^{3/4}} \tag{61}$$

Note that $q^- \gg |q^+|$.

Let us define the four-vector p'_i by the condition

$$p_i'^2 = m_i^2, \quad \frac{p'_i + P_i}{|p'_i + P_i|} = G' \tag{62}$$

Then a direct calculation shows that if P_i is fixed then [20]

$$d\rho(\mathbf{p}'_{i\perp}, p_i'^+) = \frac{|p'_i + P_i|^4}{(p'_i, p'_i + P_i)} d\rho(\mathbf{G}'_\perp, G'^+) \tag{63}$$

Therefore, as follows from Eqs. (33), (35) and (38), (39), (59), (60) and (62),

$$W^{\mu\nu} = \frac{1}{4\pi} \sum_{\sigma_1 \ldots \sigma_N} \sum_{i,j=1}^{N} \sum_{\sigma_i', \sigma_j'} \int \{ \prod_{l=1}^{N} d\rho(\mathbf{p}_{l\perp}, p_l^+) \} \frac{r_i r_j}{m^2} \cdot$$
$$(2\pi)^4 \delta^{(4)}(P' + q - P) \frac{M_0(\mathbf{k}_i', M_i)^3 M_0(\mathbf{k}_j', M_j)^3}{\omega_i(\mathbf{k}_i')\omega_j(\mathbf{k}_j')} \cdot$$
$$\langle \chi^P(\mathbf{k}_j', int_j, \sigma_1, \ldots\sigma_j', \ldots\sigma_N) | [\bar{w}_j(p_j', \sigma_j')\gamma^\mu w_j(p_j, \sigma_j)] \cdot$$
$$[\bar{w}_i(p_i, \sigma_i)\gamma^\nu w_i(p_i', \sigma_i')] | \chi^P(\mathbf{k}_i', int_i, \sigma_1, \ldots\sigma_i', \ldots\sigma_N) \rangle \tag{64}$$

where (compare with Eq. (42))

$$k_i' = L[\beta(\mathbf{G}_\perp', G'^+)]^{-1} p_i', \quad \xi_i' = \frac{\omega_i(\mathbf{k}_i') + k_i'^z}{M_0(\mathbf{k}_i', M_i)} \tag{65}$$

and $k_i' = (\omega_i(\mathbf{k}_i'), \mathbf{k}_i')$.

We can write P_i on the one hand as a function of G', p_i', int_i and on the other hand as a function of G, p_i, int_i. Therefore (compare with Eq. (43)) $\mathbf{k}_i'$ can be defined by the condition

$$L[\beta(\mathbf{G}_\perp', G'^+)]((M_i^2 + \mathbf{k}_i'^2)^{1/2}, -\mathbf{k}_i') =$$
$$L[\beta(\mathbf{G}_\perp, G^+)]((M_i^2 + \mathbf{k}_i^2)^{1/2}, -\mathbf{k}_i) \tag{66}$$

As follows from Eq. (34), in the reference frame under consideration

$$\mathbf{k}_{i\perp}' = \mathbf{k}_{i\perp}, \quad k_i'^z = (1 + 2|G^z|^2) k_i^z - 2G^0 G^z (M_i^2 + \mathbf{k}_i^2)^{1/2} \tag{67}$$

As follows from this expression, in the Bjorken limit

$$k_i'^z = (1 + 2|G^z|^2)[k_i^z + (M_i^2 + \mathbf{k}_i^2)^{1/2}] - \frac{(M_i^2 + \mathbf{k}_i^2)^{1/2}}{4|G^z|^2} \tag{68}$$

Therefore taking into account Eq. (26) we conclude that $|k_i'^z| \leq m_0$ (see Sect. 4) if $|k_i^z| \approx M/2$ and $k_i^z < 0$. Then, as follows from Eqs. (54), (67) and (68)

$$k_i'^z = \frac{M_i^2 + \mathbf{k}_{i\perp}^2}{2m(1-x)} - \frac{m(1-x)}{2} \tag{69}$$

As follows from Eqs. (66) and (69),

$$M_0(\mathbf{k}_{i\perp},\xi_i',M_i)(1-\xi_i')=m(1-x) \tag{70}$$

where the function $M_0(\mathbf{k}_{i\perp},\xi_i',M_i)$ is defined by Eq. (26).

As follows from Eqs. (26) and (70), the explicit expression of ξ_i' as a function of $\mathbf{k}_{i\perp},M_i,x$ is

$$\xi_i'=\left\{\begin{array}{l}\frac{1}{2}\{1-\alpha_i-\beta_i+[(1-\alpha_i-\beta_i)^2+4\alpha_i]^{1/2}\}\\ \frac{1}{2}\{1-\alpha_i-\beta_i-[(1-\alpha_i-\beta_i)^2+4\alpha_i]^{1/2}\}\\ \frac{m_i^2+\mathbf{k}_{i\perp}^2}{m_i^2+\mathbf{k}_{i\perp}^2+M'^2(1-x)^2}\end{array}\right. \tag{71}$$

if $M_i>m_i$, $M_i<m_i$ and $M_i=m_i$ respectively where

$$\alpha_i=\frac{m_i^2+\mathbf{k}_{i\perp}^2}{M_i^2-m_i^2},\quad \beta_i=\frac{M'^2(1-x)^2}{M_i^2-m_i^2} \tag{72}$$

Since $x\in[0,1]$, it follows from Eqs. (71) and (72) that $\xi_i'\in[\xi_i^{min},1]$ where $\xi_i^{min}=\xi_i^{min}(\mathbf{k}_{i\perp},M_i)$ is a function of $\mathbf{k}_{i\perp},M_i$ which can be determined from Eqs. (71) and (72) at $x=0$. It is easy to see that $0<\xi_i^{min}<1$.

By analogy with the consideration in the preceding section we can show that in our model the quarks absorb the virtual photon incoherently. Therefore everything is ready for the calculation of the hadronic tensor in Eq. (64). Since we wish to compare the results with those obtained in the front form, we note that, as follows from Eqs. (25), (27) and (28), $\chi^P=\chi^F/M_0$. Then a simple calculation using Eqs. (17), (24), (28), (34), (47), (54), (65), (67) and (70) shows that in the reference frame under consideration

$$\begin{aligned}W^{\mu\nu} &= \sum_{i=1}^{N} r_i^2\int\langle\chi^F(\mathbf{k}_{i\perp},\xi_i',int_i)|S_i^{\mu\nu}|\chi^F(\mathbf{k}_{i\perp},\xi_i',int_i)\rangle\cdot\\ &\quad (\frac{1-x}{1-\xi_i'})^3[1+\frac{k_i'^z}{\omega_i(\mathbf{k}_i')}]^2\frac{d^2\mathbf{k}_{i\perp}d\rho^F(int_i)}{4(2\pi)^3\xi_i'(1-x)}\end{aligned} \tag{73}$$

where $\xi_i'=\xi_i'(\mathbf{k}_{i\perp},M_i,x)$ is given by Eqs. (71) and (72).

Let us introduce the function

$$\begin{aligned}\tilde{\rho}_i(x) &= \sum_{\sigma_1...\sigma_N}\int|\chi^F(\mathbf{k}_{i\perp},\xi_i',int_i,\sigma_1,...\sigma_N)|^2\cdot\\ &\quad (\frac{1-x}{1-\xi_i'})^3[1+\frac{k_i'^z}{\omega_i(\mathbf{k}_i')}]^2\frac{d^2\mathbf{k}_{i\perp}d\rho^F(int_i)}{2(2\pi)^3\xi_i(1-x)}\end{aligned} \tag{74}$$

In contrast with the function $\rho_i(x)$ defined by Eq. (49), the function $\tilde{\rho}_i(x)$ obviously has no probabilistic interpretation. As follows from Eqs. (50), (61) and (73), in our model Bjorken Scaling and the Callan-Gross relation [37] also take place since

$$F_1(x) = \frac{1}{2}\sum_{i=1}^{N} r_i^2 \tilde{\rho}_i(x), \quad F_2(x) = 2xF_1(x) \tag{75}$$

Analogously we can calculate the hadronic tensor for the neutrino (antineutrino) - nucleon scattering, and, instead of Eq. (53), the result is

$$\begin{aligned}
&F_1^{\nu p}(x) = \sum_{i=d,s,\bar{u},\bar{c}} \tilde{\rho}_i(x), \quad F_1^{\bar{\nu}p}(x) = \sum_{i=u,c,\bar{d},\bar{s}} \tilde{\rho}_i(x),\\
&F_3^{\nu p}(x) = 2\sum_{i=d,s} \tilde{\rho}_i(x) - 2\sum_{i=\bar{u},\bar{c}} \tilde{\rho}_i(x),\\
&F_3^{\bar{\nu}p}(x) = 2\sum_{i=u,c} \tilde{\rho}_i(x) - 2\sum_{i=\bar{d},\bar{s}} \tilde{\rho}_i(x),\\
&F_2^{\nu p}(x) = 2xF_1^{\nu p}(x), \quad F_2^{\bar{\nu}p}(x) = 2xF_1^{\bar{\nu}p}(x)
\end{aligned} \tag{76}$$

In connection with the discussion of the role of off-shellness in Sect. 2, let us note that in our approach the four-momenta of interacting particles are always on-mass shell (by analogy with the "old-fashioned" time ordered perturbation theory) while in the Feynman diagram approach such four-momenta $\tilde{p}_i$ can be off-shell, i.e. in the general case $\tilde{p}_i^2 \neq m_i^2$. In diagrams which in the conventional approach are supposed to be dominant ("handbag diagrams") the four-momentum of the struck quark in the initial state is often written in the form $\tilde{p}_i = x_i P' + \tilde{k}_i$ where all the components of the vector $\tilde{k}_i$ are much smaller than Q. Therefore the four-momentum of this quark in the final state is equal to $\tilde{p}_i + q$ and the condition that the off-shellness is much smaller than Q^2 implies that $|\tilde{p}_i^2|, |(\tilde{p}_i + q)^2| \ll Q^2$. Hence we easily derive that $x_i = x$ in the Bjorken limit. It is obvious that $x_i = \tilde{p}_i{}^+/P'^+$ in the Bjorken limit and this is the argument that the Bjorken variable x can be interpreted as the momentum fraction in the IMF.

It is clear from Eq. (70) that the equality $\xi_i' = x$ takes place only if one neglects the difference between the free mass and the mass of the bound state while in the general case $\xi_i' \neq x$. In the point form ξ_i' is not equal to the quantity x_i since (see Eq. (22)) the quantity P'^+ in the point form is

not the same as for free particles. Therefore the fact that $\xi_i' \neq x$ does not contradict the relation $|\tilde{p}_i^2| \ll Q^2$. Let us calculate the quantity $\tilde{p}_i^2 - m_i^2$ in the parton model and in our one. In both cases $\tilde{p}_i = P' - P_i$ and we take into account that the vector P_i is on-shell, i.e. $P_i^2 = M_i^2$. As follows from Eqs. (46) and (67), in both cases the $\perp$ components of the vector P_i is equal to $-\mathbf{k}_{i\perp}$. Therefore in both cases $P_i^- = (M_i^2 + \mathbf{k}_{i\perp}^2)/2P_i^+$. In the parton model P_i is given by the left-hand-side of Eq. (43), and, as follows from Eqs. (42), (46) and the condition $\xi_i' = x$, $P_i^+ = (1-x)P'^+$. In our model P_i is given by the left-hand-side of Eq. (66), and, as follows from the conditions $G' = P'/m$ and Eqs. (65), (67) and (70), again $P_i^+ = (1-x)P'^+$. Therefore a simple calculation using Eq. (26) shows that in both cases

$$\tilde{p}_i^2 - m_i^2 = x[m^2 - M_0(\mathbf{k}_{i\perp}, x, M_i)^2] \tag{77}$$

On the other hand, let us calculate the quantity $q_i = p_i - p_i'$ which has the meaning of the four-momentum transferred to the struck quark. Since $P' + q = P$ and $P = P_i + p_i$, we have $q_i = q + P' - P_i - p_i'$. Since in the front form the $+, \perp$ components of P' are the same as for noninteracting particles it is obvious that $q_i^+ = q^+$ and $\mathbf{q}_{i\perp} = \mathbf{q}_\perp$. At the same time, in the reference frames considered in the preceding and this sections the minus components of P', P_i and p_i' are infinitely small. Therefore, in the parton model $q_i = q$ and in our one $q_i^- = q^-$. Therefore, as follows from Eqs. (17), (54), (61), (67) and (70), in our model

$$\mathbf{q}_{i\perp} = 0, \quad q_i^+ = q^+ \frac{(1-x)\xi_i'}{(1-\xi_i')x}, \quad q_i^- = q^-, \quad q_i^2 = q^2 \frac{(1-x)\xi_i'}{(1-\xi_i')x} \tag{78}$$

In the conventional theory we must have $q_i = q$ since the inequality of these quantities is prohibited by the factorization theorem [13, 14, 15, 16] which (see the discussion in Sect. 2) asserts in particular that upper parts of Feynman diagrams dominating in DIS are fully defined by the hard part of the DIS process. Roughly speaking this implies that the vertex describing absorption of the virtual photon by the struck constituent is not surrounded by virtual particles and therefore it is obvious that $q_i = q$. However, as argued in Sect. 2, the factorization theorem does not take place in the general case. Since in relativistic models the values of M_0 essentially differ from m, it is obvious from Eq. (77) that the off-shellness is important. For this reason it is not clear why the vertex describing the absorption of the virtual photon can be considered in perturbation theory.

Let us now consider a question which seems to be important for understanding the difference between our model and the parton one. In both cases the nucleon is described by the wave function $\varphi(p_1, ...p_N)$ where the spin variables are dropped. The quantity p_i is the four-momentum of *free* particle i. The question arises whether p_i can be interpreted as the four-momentum of particle i inside the nucleon when the particles interact with each other. As pointed out by Coester [43], such an interpretation implies that

$$[\hat{P}, p_i] = 0, \quad \hat{U}(l)^{-1} p_i^{\mu} \hat{U}(l) = L(l)^{\mu}_{\nu} p_i^{\nu} \tag{79}$$

where $\hat{U}(l)$ is the representation operator corresponding to $l \in SL(2, C)$. It is sufficient to satisfy these relations not on the whole Hilbert space but only on the subspace H_0 consisting of all possible single-nucleon states. The second relation is obviously valid only in the point form since in the front one the operators $\hat{U}(l)$ are generally speaking interaction dependent. The first relation in the point form is valid too, since, as follows from Eq. (32), the operator $\hat{P}$ in H_0 is the operator of multiplication by mG. We conclude that p_i has the meaning of the four-momentum of particle i inside the nucleon only in the point form.

Let P_i' be the initial four momentum of the rest of the target when the virtual photon is absorbed by particle i and P_i be the same quantity in the final state. While the initial state is the nucleon with the four-momentum P', the final state is the state of free particles with the four-momentum P. If the virtual photon with the four-momentum q is absorbed by particle i, then IA for the operator $\hat{J}^{\mu}(0)$ implies that the four-momentum of the spectator does not change, i.e.

$$P_i' = P_i, \quad P = P' + q \tag{80}$$

where the second relation follows from the total four-momentum conservation. It is easy to see that Eq. (43) is a consequence of Eq. (80) in the front form while Eq. (66) is a consequence of Eq. (80) in the point form. As shown in the preceding section, the parton model result $\xi_i' = x$ follows from Eq. (43) while, as shown in this section, the relation between ξ_i' and x given by Eq. (70) follows from Eq. (66). Taking into account the discussion in the preceding paragraph we again conclude that IA for the operator $\hat{J}^{\mu}(0)$ has the physical meaning only in the point form, and the derivation of the relation $\xi_i' = x$ in the parton model is not substantiated. Our discussion also gives grounds to think that, although the expression (73) for the hadronic

tensor is essentially model dependent, the relation (70) is in fact only kinematical since it does not depend on any specific features of the rest of the target (in particular on whether N is finite or infinite). Therefore, if we assume that the final state interaction is not important in the Bjorken limit, then the relation (70) is rather general, but in the general case M_i entering into the expression for M_0 is the physical (rather than the free) mass of the rest of the target.

6 Deviation from the standard sum rules

If we assume that Eq. (11) is valid beyond perturbation theory (although, as noted in Sect. 2, there are serious reasons to doubt whether this is the case) but no form of the operators $\hat{O}_i^{\mu_1\cdots\mu_n}$ is prescribed then all standard results about the Q^2 evolution of the structure functions remain. Indeed in this case the only information about the operators $\hat{O}_i^{\mu_1\cdots\mu_n}$ we need is their tensor structure since we should correctly parametrize the matrix elements $\langle P',\chi|\hat{O}_i^{\mu_1\cdots\mu_n}|P',\chi\rangle$.

However the derivation of sum rules in DIS requires additional assumptions. It is well-known that they are derived with different extent of rigor. For example, the Gottfried and Ellis-Jaffe sum rules [44, 45] are essentially based on model assumptions, the sum rule [46] was originally derived in the framework of current algebra for the time component of the current operator while the sum rules [47, 48, 49] also involve the space components. Therefore in the framework of current algebra the sum rule [46] is substantiated in greater extent than the sum rules [47, 48, 49] (for a detailed discussion see Refs. [50, 6, 8]). Moreover, there exist field theory models where $\hat{J}^0(\mathbf{x})$ is free while $\hat{\mathbf{J}}(\mathbf{x})$ is necessarily interaction dependent (see e.g. the calculations in scalar QED in Ref. [51]). Now the sum rules [46, 47, 48, 49] are usually considered in the framework of the OPE and they have the status of fundamental relations which in fact unambiguously follow from QCD. However the important assumption in deriving the sum rules is that the expression for $\hat{O}_V^\mu$ coincides with $\hat{J}^\mu$, the expression for $\hat{O}_A^\mu$ coincides with the axial current operator $\hat{J}_A^\mu$ etc. (see Eq. (12)). Our consideration in Sect. 2 shows that this assumption has no physical ground. Therefore although (for some reasons) there may exist sum rules which are satisfied with a good accuracy, the statement that the sum rules [46, 47, 48, 49] unambiguously follow from QCD is not substantiated.

As already noted, when Q^2 is large, the results of the present theory agree with the parton model up to anomalous dimensions and perturbative QCD corrections. The sum rules [46, 47, 48, 49] have the property that the corresponding anomalous dimensions are equal to zero. Therefore, if Q^2 is so large that $\alpha_s(Q^2)$ is small, the sum rules [46, 47, 48, 49] are in agreement with the parton model.

The existing experimental data do not make it possible to verify the sum rules [46, 48] with a good accuracy, and rather precise data exist for the sum rules [47, 49].

The data recently obtained by the CCFR collaboration [52] show that the experimental value of the Gross-Llewellyn Smith sum S_{GLS} is smaller than the value $S_{GLS} = 6$ predicted by the parton model. The analysis of the CCFR data in papers [52, 53, 54] shows that actually $S_{GLS} = 4.90 \pm 0.16 \pm 0.16$ at $Q^2 = 3 \cdot GeV^2$ which is smaller than the value for S_{GLS} which follows from the present theory, even if the corrections of order $\alpha_s(Q^2)$ and $\alpha_s(Q^2)^2$ are taken into account.

The data on the Bjorken sum rule [47] have been analyzed by several authors (see e.g. Refs. [55, 56, 57, 58, 59, 60, 61, 62]). While some authors argue that the data on the Bjorken sum S_B are in agreement with the theoretical prediction 0.20, another authors (see e.g. Ref. [62]) argue that the value consistent with the data of three experimental groups is 0.15.

It is important to note that the data on the sum rules [47, 49] at small values of the Bjorken variable x are known only at rather small values of Q^2 and the higher twist corrections are essentially model dependent. The problem also exists whether the extrapolation of these data to low x is correct (see the discussion in Refs. [58, 59, 60, 62]). In addition, the problem exists whether it is possible to extract the neutron structure functions from the proton and deuteron data [63].

Let us now consider the sum rules which are not considered fundamental and one way or another are based on the parton model. The present data make it possible to unambiguously conclude that the sum rules [44, 45] are not satisfied. Namely, according to the recent precise results of the NMC collaboration [64] the experimental value of the integral defining the Gottfried sum rule [44] is equal to 0.235 ± 0.026 instead of $1/3$ in the parton model, and the EMC result [65] for the first moment of the proton polarized structure function $g_1(x)$ is $\Gamma_p = 0.126 \pm 0.010 \pm 0.015$ while the Ellis-Jaffe sum rule

[45] predicts $\Gamma_p^{EJ} = 0.171 \pm 0.004$.

The most impressive results of the parton model are those concerning the quark contribution to the nucleon momentum and spin.

The first result (see, for example, the discussions in Ref. [66]) says that quarks carry only 46% of the nucleon momentum, and this fact is usually considered as one of those which demonstrates the existence of gluons. In the framework of the OPE the gluon contribution to the nucleon momentum also can be calculated (see e.g. Refs. [67, 68, 12]) and the result is in agreement with the data. However the theoretical formulas are not well substantiated since corrections to them have the form $K\alpha_s(Q^2)^{-d}$ where $d > 0$ and the coefficient K is not known [68, 12].

The second result known as "the spin crisis" says that the quark contribution to the nucleon spin is comparable with zero (a detailed discussion of the spin crisis can be found, for example, in Refs. [55, 58, 56, 57, 60]).

Of course, these results are not in direct contradiction with constituent quark models since the latter are successful only at low energies. Nevertheless, our experience can be hardly reconciled with the fact that the role of gluons is so high.

The above discussion gives grounds to conclude that in the parton model the values given by the sum rules systematically exceed the corresponding experimental quantities while the quark contribution to the nucleon momentum and spin is underestimated.

The standard sum rules use the fact that in the parton model $\xi_i' = x$, and the functions $\rho_i(x)$ (see Eqs. (49)) satisfy the normalization condition

$$\int_0^1 \rho_i(x)dx = 1 \tag{81}$$

However in our model the structure functions depend on the functions $\tilde{\rho}_i(x)$ (see Eqs. (74-76)) which do not satisfy such a condition. The matter is that, as follows from Eqs. (71) and (72), the DIS data do not make it possible to determine the quark distribution at $\xi_i' < \xi_i^{min}$. The normalization integrals contain the integration over $\xi_i' \in [0, 1]$ while the DIS data make it possible to determine some integrals over $x \in [0, 1]$.

Using Eqs. (22), (26), (70), (74) and changing the integration variable from x to ξ_i' we get

$$\tilde{\rho}_i \equiv \int_0^1 \tilde{\rho}_i(x)dx = \sum_{\sigma_1,\ldots\sigma_N} \int \frac{d^2\mathbf{k}_{i\perp} d\rho^F(int_i)}{(2\pi)^3} \int_{\xi_i^{min}}^1 \frac{d\xi_i'}{2\xi_i'(1-\xi_i')}\,.$$

$$(\frac{1-x}{1-\xi_i'})^3|\chi^F(\mathbf{k}_{i\perp},\xi_i,int_i,\sigma_1,...\sigma_N)|^2[1+\frac{k_i'^z}{\omega_i(\mathbf{k}_i')}] \tag{82}$$

As argued in the preceding section, this expression is essentially model dependent but the relation (70) is rather general. In particular ξ_i^{min} can considerably differ from zero. Therefore, comparing Eq. (82) with Eqs. (49) and (81), it is natural to expect that $\tilde{\rho}_i < 1$.

Let us consider, for example, the Gottfried sum rule [44], according to which the quantity

$$S_G = \int [F_{2p}(x) - F_{2n}(x)]\frac{dx}{x} \tag{83}$$

is equal to 1/3. Here $F_{2p}(x)$ and $F_{2n}(x)$ are the structure functions F_2 for the proton and neutron respectively. This sum rule easily follows from Eqs. (49), (51), (81) if we assume that the neutron wave function can be obtained from the proton one if one of the u quarks in the proton is replaced by the d quark. We suppose that particle 1 in the proton is the u quark, particle 1 in the neutron is the d quark and all other particles are the same. Then, as follows from Eqs. (75), (82) and (83), $S_G = \tilde{\rho}_1/3$. Therefore, if $\tilde{\rho}_1 < 1$ then $S_G < 1/3$.

The Gross-Llewellyn Smith sum rule [49] reads

$$S_{GLS} = \int_0^1 [F_3^{\bar{\nu}p}(x) + F_3^{\nu p}(x)]dx = 6[1 + (...)] \tag{84}$$

where (...) stands for the terms of order $\alpha_s(Q^2)$, $\alpha_s(Q^2)^2$ etc. It is obvious from Eqs. (53) and (81) that the results of the parton model agree with Eq. (84) in leading order in $\alpha_s(Q^2)$. However, our result, which follows from Eqs. (76) and (82), is that if, for simplicity, we assume that the values of $\tilde{\rho}_i$ for the sea quarks and the corresponding antiquarks are the same, then in leading order in $\alpha_s(Q^2)$

$$S_{GLS} = 4\tilde{\rho}_u + 2\tilde{\rho}_d \tag{85}$$

where $\tilde{\rho}_u$ and $\tilde{\rho}_d$ are the quantities $\tilde{\rho}_i$ for the valence u and d quarks respectively. Since it is natural to expect that $\tilde{\rho}_u, \tilde{\rho}_d < 1$, one also might expect that $S_{GLS} < 6$.

Analogously, the DIS data alone do not make it possible to determine the contributions of the u, d and s quarks to the nucleon spin, and the problem of the spin crisis does not arise. Indeed, these contributions (usually denoted

as $\Delta q = (\Delta u, \Delta d, \Delta s))$ are given by some integrals over $\xi_i \in [0,1]$. Since the integrals over x can be transformed to the integrals over $\xi_i \in [\xi_i^{min}, 1]$, we see that the DIS data do not make it possible to determine the contributions of $\xi_i \in [0, \xi_i^{min}]$ to Δq. Thus it is natural to expect that the parton model underestimates the quantities Δq.

As noted by many authors, the Gottfried sum rule is not a consequence of (perturbative) QCD. Therefore it is not strange that this sum rule disagrees with the experimental data, and different explanations of the disagreement were proposed. Analogously, there exist many papers devoted to the Ellis-Jaffe sum rule and to the spin crisis. However there are no approaches explaining why the parton model overestimates the experimental quantities for all the sum rules and underestimates the quark contribution to the nucleon momentum and spin. These facts are qualitatively explained in our model.

7 Discussion

According to our experience in conventional nuclear and atomic physics, in processes with high momentum transfer the effect of binding is not important and the current operator can be taken in IA, i.e. this operator is the same as for the system of free particles. However this experience is based on the nonrelativistic quantum mechanics where only the Hamiltonian is interaction dependent and the other nine representation generators of the Galilei group are free, while in the relativistic case at least three representation generators of the Poincare group are interaction dependent (see Sect. 1). The usual motivation of IA is that if Q is very large then the process of absorption of the virtual photon is so quick that the struck quark does not interact with other constituents during this process. Such a motivation is reasonable in the nonrelativistic case where the kinetic energies and interaction operators in question are much smaller than the masses of the constituents. However in the relativistic case the off-shellness of the struck quark is important and the properties of the off-shell quark are not the same as of the free one.

The parton model is equivalent to IA in the front form of dynamics (see the discussion in Sect. 4). The usual motivation of the parton model is that, the partons in the IMF are free to the extent that IA is valid though the partons form the bound state—the nucleon. It is stated that such a duality is a consequence of asymptotic freedom. In our opinion, this argument is

not convincing since asymptotic freedom is the argument in favor of using perturbation theory in $\alpha_s(Q^2)$ but the bound state cannot be considered in such a framework.

In the present theory of DIS the parton model is a consequence of the factorization theorem [13, 14, 15, 16] and the OPE [17]. It is well-known that the OPE has been proved only in perturbation theory but many physicists believe that there can be no doubt about the validity of the OPE beyond that theory. However the consideration in Sect. 2 shows the OPE beyond perturbation theory is problematic and in the general case the factorization theorem does not work. Therefore the current operator contains a nontrivial nonperturbative part which contributes to DIS even in the Bjorken limit. Since this contribution cannot be determined at the present stage of QCD, it is reasonable to consider models in which the representation operators of the extended Poincare and the current operator can be explicitly constructed and it is possible to verify that they satisfy proper commutation relations.

The important fact (which in our opinion has been overlooked by physicists working on DIS) is that *in the parton model the conditions (6) and (36) are not satisfied and therefore Lorentz invariance, P invariance and T invariance of the current operator are violated*. This fact has been explained in detail in Sects. 2, 4 and 5. What is the extent of the violation of these symmetries in the parton model? The results obtained in the parton model are the same as in the OPE if the anomalous dimensions are discarded and only the leading terms of perturbation theory are taken into account. Since the anomalous dimensions can be determined in the framework of perturbative QCD, the difference between the current operators in the OPE and in the parton model also can be determined in such a framework. Meanwhile the consideration in Sects. 4 and 5 shows that the reason of the violation of extended Lorentz invariance in the parton model is the difference between the real mass operator $\hat{M}$ and the free mass operator M_0. This difference is responsible for binding of quarks and gluons in the nucleon and it cannot be determined in the framework of perturbative QCD. This observation is the additional argument that perturbative QCD does not apply to DIS.

In Sect. 5 we consider a model in which IA for the operator $\hat{J}^\mu(0)$ is compatible with extended Poincare invariance and current conservation. While the model cannot have any pretensions to be fundamental, we argue that the relation (70) between the light cone momentum fraction ξ_i' and the Bjorken variable x is in fact kinematical. Therefore DIS experiments alone cannot

determine the ξ_i' distribution of quarks in the nucleon. This conclusion also poses the problem whether the parton densities extracted from the data on DIS at the assumption $\xi_i' = x$ can be used for describing Drell-Yan processes, hard $p\bar{p}$ collisions etc.

If the bound state under consideration can be described nonrelativistically then the results of the parton model and our one are the same since in the nonrelativistic approximation $M_0 \approx m \approx m_1 + ...m_N$. However there is no reason to think that the nucleon is the nonrelativistic system of quarks and gluons.

The arguments that nonperturbative effects in the current operator are important even in the Bjorken limit were given by several authors (see e.g. Refs. [69, 70, 71, 72, 73, 74, 75]). The main difference between our consideration and the consideration in these references is that we consider in detail the restrictions imposed on the current operator by its commutation relations with the representation operators of the extended Poincare group. In the conventional theory of DIS as well it has never been verified that the factorization theorem and the OPE beyond perturbation theory are compatible with these restrictions.

Although there is no doubt that predictions of perturbative QCD are in qualitative agreement with experimental data, the extent to which the agreement is quantitative is not quite clear (we do not discuss this question in the present paper). Therefore it is very important to carry out experiments the results of which will undoubtedly show wether perturbative QCD applies to DIS. In particular it is very important to test the sum rules [46, 47, 48, 49]. However for the unambiguous test they should be extracted directly from the experimental data at large Q^2, not using the Q^2 evolution determined from the OPE or Altarelli-Parisi equations [76]. As argued in Ref. [63], it is also not clear whether the neutron structure functions at small x can be extracted from the proton and deuteron data even in principle. In this reference we also argue that the experiment which will show whether the factorization property takes place, is deuteron DIS at large Q^2 and small x.

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